From Coulomb blockade to the Kondo regime in a Rashba dot

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We investigate the electronic transport in a quantum wire with localized Rashba interaction. The Rashba field forms quasi-bound states which couple to the continuum states with an opposite spin direction. The presence of this Rashba dot causes Fano-like antiresonances and dips in the wire's linear conductance. The Fano lineshape arises from the interference between the direct transmission channel along the wire and the hopping through the Rashba dot. Due to the confinement, we predict the observation of large charging energies in the local Rashba region which lead to Coulomb-blockade effects in the transport properties of the wire. Importantly, the Kondo regime can be achieved with a proper tuning of the Rashba interaction, giving rise to an oscillating linear conductance for a fixed occupation of the Rashba dot.

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I. INTRODUCTION

The high degree of functionality of spin-based electronic devices has attracted much attention due to promising applications in quantum and classical computation.[1] Use of spins to store and carry classical information has been proven to be both faster and less power consuming than the conventional technology based on the control of the charge. Besides, spin-1/2 systems are genuine two-level systems and, therefore, natural building blocks for quantum computation.[2] These features have motivated the emergence of an exciting area of research termed *spintronics*.

Two-dimensional (2D) semiconductors are appropriate materials to accomplish spintronic applications since they offer the possibility of electric control of spins via tunable *spin-orbit* interactions. The tuning is achieved with electric fields induced by external gates coupled to the semiconductor.[3] A prominent contribution to spin-orbit effects in 2D electron gases of narrow-gap semiconductors (typically, InAs materials) is the Rashba interaction.[4] This is caused by an asymmetry in the potential defining the quantum well in the direction perpendicular to the 2D electron gas. The control of the Rashba coupling strength opens the possibility of investigating 2D systems with spatially modulated Rashba fields such as those having constant Rashba strength in a semiplane[5], a stripe[6] or an island.[7,8,9,10,11] Less known are the effects of a local Rashba field on the electron-electron interaction. Our aim is to examine how sensitive the charging energy of a quasi-1D Rashba region is to changes of the spin-orbit strength and the consequences in the transport properties.

The Rashba Hamiltonian reads

$$\mathcal{H}_{\rm R} = \frac{1}{2\hbar}(\{\alpha, p_y\}\sigma_x - \{\alpha, p_x\}\sigma_y)\,, \qquad (1)$$

where α is the strength while $\vec{p}$ and $\sigma_{x,y,z}$ are the 2D momentum operator and the Pauli matrices, respectively. The anticonmutators in Eq. (1) ensure a Hermitian $\mathcal{H}_{\rm R}$ when α is spatially nonuniform. For strictly 1D systems the precession term in Eq. (1), $\{\alpha, p_x\}\sigma_y$ for transport along x, leads to the formation of bound states.[11] These bound states acquire a finite lifetime in quasi-1D systems where the term $\{\alpha, p_y\}\sigma_x$ couples adjacent subbands with opposite spin directions.[12] In a quantum wire these subbands arise from the parabolic confinement of the 2D gas. Therefore, a local Rashba interaction in a quantum wire has two main effects, namely *(i) it forms bound states* in each subband and *(ii) it broadens these bound states due to coupling with adjacent subbands*. Hereafter, we refer to the quasibound states as *Rashba dots*.[11]

In this work we are interested in the transport properties of a quantum wire with a local Rashba interaction in the presence of Coulomb repulsion. As we have partially anticipated, the physical scenario is governed by the interference between a direct (nonresonant) transmission channel with a path that passes through the Rashba dot, leading to the formation of *Fano* resonances[11,13] as shown in Fig. 1(a). Due to the confinement of the electron motion within the Rashba dot our calculations show that the Coulomb interaction becomes in fact very large, see Fig. 1(b). Thus, it is possible to observe charging effects where the transport through the Rashba dot is governed by Coulomb blockade. Furthermore, a proper combination of gates leads to a very strong coupling of a singly occupied Rashba dot to the continuum states. In this case and when the Coulomb interaction is sufficiently large the *Kondo effect*[14,15] takes place at very low temperatures. Then, the localized spin in the Rashba dot forms a many-body singlet state with the continuum electron spins and, hence, is screened. This occurs at energies k_BT lower than the Kondo energy scale k_BT_K, which is the binding energy of the many-body singlet state.[14,15] It manifests itself as a quasi-particle resonance, namely, a

peak of width $k_B T_K$ of the local density of states (LDOS) at the Fermi energy (E_F). Experiments have been able to observe Kondo effects in quantum dots,[16] its most remarkable signature being the *unitary limit* of the linear conductance at $k_B T = 0$, i.e., $\mathcal{G}_0 = 2e^2/\hbar$.[15] In our case, the Kondo resonance that forms in the Rashba dot destructively interferes with the nonresonant transmitting mode giving rise to an oscillatory $\mathcal{G}_0$.

This work considers the transport properties of a quantum wire with a localized Rashba region including charging and correlation effects. Section II presents the model of Hamiltonian developed to describe this system. Section III is devoted to investigate the linear conductance across the quantum wire with the localized Rashba interaction. The results derived from our calculations are shown in Section III. Finally, the main conclusions of our work are summarized in Section IV.

II. MODEL

In this section we discuss the model Hamiltonian that describes a quantum wire with a localized Rashba interaction. As emphasized, the Rashba interaction plays the role of both (i) the attractive potential and (ii) the coupling to the continuum states. Each subband couples with at least one bound state (Rashba dot) splitting off the higher subband. For simplicity we consider an energy range in the first plateau, $\varepsilon_1 < E < \varepsilon_2$ with $\varepsilon_\nu = (\nu - 1/2)\hbar\omega_0$ being the energy of the transversal modes (ω_0 defines the confinement strength of the quasi-1D system). The bound state energy $\epsilon_0 = \varepsilon_2 + \epsilon_b$ lies within the first plateau because of its negative binding energy ϵ_b. Thus, the first subband consists of a continuum of states and the second subband provides the bound state which is coupled to the first subband states with opposite spin. As a consequence, the transport can occur through two different paths, (i) via the quasi-bound state (through the Rashba dot described by an Anderson-like Hamiltonian) or (ii) via a nonresonant path through the continuum states of the first subband. The Hamiltonian reads:

$$\mathcal{H} = \sum_{\alpha k\sigma} \epsilon_{\alpha k}\, c^\dagger_{\alpha k\sigma} c_{\alpha k\sigma} + \sum_{k\sigma} (W e^{i s_\sigma \varphi} c^\dagger_{Lk\sigma} c_{Rk\sigma} + \text{H.c.}) + \sum_\sigma \epsilon_0\, d^\dagger_\sigma d_\sigma + U n_\sigma n_{\bar\sigma} + \sum_{\alpha k\sigma} (V_0 c^\dagger_{\alpha k\sigma} d_{\bar\sigma} + \text{H.c.})\,. \quad (2)$$

We choose the spin quantization axis along the Rashba field (the y-axis for propagation along x). $d^\dagger_\sigma$ creates an electron with spin $\sigma = \uparrow, \downarrow$ in the Rashba dot ($n_\sigma = d^\dagger_\sigma d_\sigma$ is the Rashba dot occupation per spin σ) and $c^\dagger_{L(R)k\sigma}$ creates a continuum electron with k wavevector and spin σ in the left (right) side of the Rashba dot. Both electronic states are coupled via the hopping amplitude V_0. We note that this hopping originates from the Rashba interaction only,[11] and that this qualitatively differs from the coupling via tunnel barriers as in conventional quantum dot models. We define the signs $s_{\uparrow,\downarrow} = \pm 1$ and $\bar\sigma$ indicates reverse spins. U is the on-site Coulomb interaction in the Rashba dot. Depending on the strength of U, which we calculate below, charging and correlation effects can be present. We note that the Hamiltonian (2) presents a great similarity with that describing the electronic transport in a closed Aharanov-Bohm (AB) interferometer with a quantum dot in one of its arms.[17,18,19,20] Nevertheless, there are two important differences. First, the phases of the hopping amplitudes W in Eq. (2) depend on the spin index. Such spin-dependent phases have been shown to give rise to spin polarization in interferometer setups.[21] Second, each hopping process through the dot is associated with a spin-flip event. Spin-flip interactions have been recently considered in strongly correlated quantum dots[22] but mainly dealing with intradot spin-flip scattering. Spin-flip assisted tunneling has received much less attention.[21,23]

In the usual Anderson Hamiltonian, spin is conserved during tunneling and this leads to Kondo correlations at low temperatures. We now show that, despite the fact that Eq. (2) involves spin-flip hopping due to the Rashba interaction, the Kondo effect also manifests itself in this system. For simplicity, let us neglect the nonresonant path in the following discussion ($W = 0$). We first perform a Schrieffer-Wolff transformation to the Rashba dot Hamiltonian. The resulting Hamiltonian,

$$\mathcal{H}'_K = J_{\alpha\beta}(S^x s^x_{\alpha\beta} - S^y s^y_{\alpha\beta} - S^z s^z_{\alpha\beta})\,, \quad (3)$$

is equivalent to the usual Kondo Hamiltonian

$$\mathcal{H}_K = \sum_{\alpha\beta} J_{\alpha\beta} \vec{S} \cdot \vec{s}_{\alpha\beta}\,, \quad (4)$$

where the spin of the Rashba dot $\vec{S}$ interacts antiferromagnetically ($J_{\alpha\beta} > 0$) with the spin density of the delocalized electrons $\vec{s}_{\alpha\beta} = \psi^\dagger_{\alpha\nu} \vec{\sigma}_{\nu\mu} \psi_{\beta\mu}$ ($\psi_{\alpha\nu} = \sum_k c_{\alpha k\nu}$ are the conduction field operators at the Rashba dot site). To see this, we can perform a unitary transformation to convert $\mathcal{H}'_K$ into $\mathcal{H}_K$, namely, we perform a rotation of π around the x axis in the localized spin space keeping the delocalized spins unchanged, i.e.,

$$S^{y(z)} \dashrightarrow -S^{y(z)}, \quad S^x \dashrightarrow S^x\,. \quad (5)$$

In this way we transform $\mathcal{H}'_K$ into $\mathcal{H}_K$.
Physically, the system we consider is characterized by just two parameters: α, the intensity of the Rashba field, and ℓ, the length where this is active in the quantum wire [see inset of Fig. 1(b)]. Importantly, the parameters governing the Hamiltonian (2), $U \equiv U(\alpha, \ell)$, $V_0 \equiv V_0(\alpha, \ell)$, $W \equiv W(\alpha, \ell)$, and $\varphi(\alpha, \ell)$ can be externally controlled by changing both α and ℓ. The Rashba phase $\varphi = k_\alpha \ell$ with $k_\alpha = m\alpha/\hbar^2$ is the total phase gained by an electron traveling from the left to the right side of the Rashba dot.

III. LINEAR CONDUCTANCE

Once we have established, in the previous section, that in a quantum wire with a localized Rashba interaction Kondo physics can arise, now we derive the transport properties for such a system.

The electrical current is derived from the time derivative of the quantum occupation of the conduction electrons on the left side of the Rashba dot:

$$I_L = ed\langle N_L\rangle/dt = (ie/\hbar)\langle[\mathcal{H}, N_L]\rangle\,, \qquad (6)$$

where $N_L = \sum_{k\sigma} c^\dagger_{Lk\sigma} c_{Lk\sigma}$. In terms of the nonequilibrium Green functions $G^<_{LR\sigma} = i\langle c^\dagger_{Rk\sigma} c_{Lk\sigma}\rangle$ and $G^<_{d\bar\sigma,L\sigma} = i\langle c^\dagger_{Lk\sigma} d_{\bar\sigma}\rangle$ the current reads

$$I_L = \frac{2e}{h}\mathcal{R}e\sum_{k\sigma}\int d\epsilon\Big[W e^{is_\sigma\varphi} G^<_{LR\sigma}(\epsilon) + V_0 G^<_{d\bar\sigma,L\sigma}(\epsilon)\Big]. \quad (7)$$

After some algebra, the current traversing the system can be written in terms of the transmission as

$$I = (e/h)\sum_\sigma \int d\epsilon \mathcal{T}_\sigma(\epsilon)[f_L(\epsilon) - f_R(\epsilon)]\,, \qquad (8)$$

where $f_{L(R)}$ is the left (right) Fermi function and the transmission reads

$$\begin{aligned}\mathcal{T}_\sigma(\epsilon) = T_b &+ 2\sqrt{T_b R_b}\cos(s_\sigma\varphi)\tilde\Gamma\mathcal{R}e G^r_{d\bar\sigma}(\epsilon)\\ &- \tilde\Gamma\left\{\left[1 - T_b\cos^2(s_\sigma\varphi)\right] - T_b\right\}\mathcal{I}m G^r_{d\bar\sigma}(\epsilon)\,. \quad (9)\end{aligned}$$

In Eq. (9),

$$T_b = 4x^2/(1+x^2)^2,\ (R_b = 1 - T_b)\,, \qquad (10)$$

is the background transmission with $x = \pi W\rho$ (ρ is the DOS of the conduction electrons), and $\tilde\Gamma = \Gamma/(1+x^2)$ with $\Gamma = \pi V_0^2\rho$. $G^r_{d\sigma}$ is the retarded Green function for the Rashba dot. This Green function is determined by calculating its self-energy which takes into account the single-particle self-energy due to tunneling and the many-body effects enclosed in the interacting self-energy $\Sigma_{\rm int}(\epsilon)$:

$$G^r_{d\sigma}(\epsilon) = \frac{1}{\epsilon - \epsilon_{0\sigma} + i\tilde\Gamma - \Sigma_{\rm int}(\epsilon)}\,. \qquad (11)$$

Notice that Eq. (9) is equivalent to the transmission obtained in Refs. 19,20 for the closed AB interferometer since

$$\cos(s_\sigma\varphi) = \cos(\varphi)\,, \quad G^r_{d\sigma} = G^r_{d\bar\sigma}\,, \qquad (12)$$

whenever the spin degeneracy is not broken. In linear response the conductance is just

$$\mathcal{G}_0 \equiv \lim_{V\to 0}\frac{dI}{dV} = \frac{2e^2}{h}\mathcal{T}(E_F)\,. \qquad (13)$$

The spin-degenerate transmission at the Fermi energy $\mathcal{T}(E_F)$ can be written in an extended Fano form as

$$\mathcal{T}(E_F) = T_b\frac{|\xi + q|^2}{\xi^2 + 1}\,, \qquad (14)$$

with a *complex Fano parameter*

$$q = \sqrt{\frac{1}{T_b}}\left(\sqrt{R_b}\cos\varphi + i\sin\varphi\right)\,, \qquad (15)$$

and

$$\xi = \frac{2\left[E_F - \epsilon_0 - \mathcal{R}e\Sigma_{\rm int}(E_F)\right]}{\tilde\Gamma - \mathcal{I}m\Sigma_{\rm int}(E_F)}\,. \qquad (16)$$

Importantly, q is complex even though the Rashba interaction is time-reversal invariant. This is in clear contrast with the Aharanov-Bohm case where the complex Fano factor is usually attributed to the broken time-reversal symmetry by the magnetic field.[24] Therefore, further investigation should be necessary to clarify the general conditions for the appearance of a complex Fano factor.[25]

A. Zero temperature: Friedel Langreth sum rule

At $k_BT = 0$ the quasiparticles in an interacting system (a Fermi liquid) possess an infinite life-time, meaning that the imaginary part of the interacting self-energy vanishes, i.e., $\mathcal{I}m\Sigma_{\rm int}(E_F) = 0$. According to this, the Friedel-Langreth sum rule[26] relates the scattering phase shift δ, picked up by the electrons when hopping, with the total quantum occupation $\langle n\rangle = \sum_\sigma\langle n_\sigma\rangle$ of the interacting system

$$\delta = \frac{\pi\langle n\rangle}{2}\,. \qquad (17)$$

The quantum occupation for the interacting system at $T = 0$ can be easily obtained in terms of ξ

$$\langle n\rangle = \frac{2}{\pi}\cot^{-1}\xi\,. \qquad (18)$$

As a consequence, the transmission at the Fermi level, $\mathcal{T}(E_F)$, can be cast in terms of δ only. Moreover, in the pure Kondo regime the dot level is renormalized by the interactions and pinned at the Fermi energy, i.e.,

$$\epsilon_0 + \mathcal{R}e\Sigma_{\rm int}(E_F) = E_F\,, \qquad (19)$$

giving rise to the Kondo resonance. As a result, the occupation per spin in the pure Kondo regime does not fluctuate and remains constant: $\langle n_\sigma\rangle = 1/2$. This fact implies that the linear conductance has a simple form:[20]

$$\mathcal{G}_0 = \frac{2e^2}{h}\left(1 - T_b\cos^2\varphi\right)\,. \qquad (20)$$

Unlike the AB interferometer where φ is constant, now the phase changes with the Rashba intensity. This means

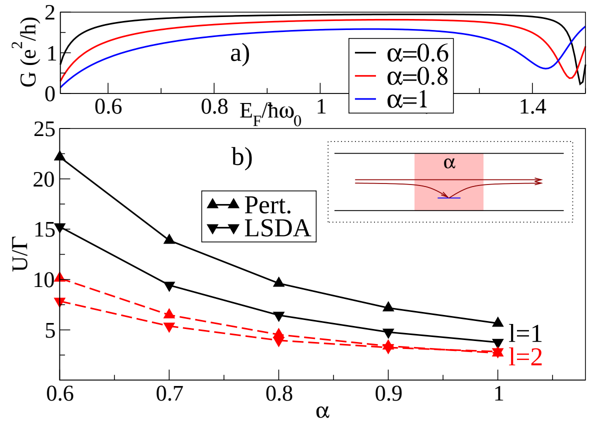

FIG. 1: (a) (Color online) (a) Linear conductance $\mathcal{G}_0$ of a noninteracting quantum wire as a function of the Fermi energy E_F (in units of $\hbar\omega_0$) with $\ell = 1$ (in units of $\ell_0 = \sqrt{\hbar/m\omega_0}$) and different Rashba strengths α (in units of $\hbar\omega_0\ell_0$). A dip forms close to the onset of the second plateau. It follows from the formation of a quasibound state which couples to the conduction subband [see Inset in (b) for a sketch], giving rise to a Fano lineshape in $\mathcal{G}_0$. Both the position and broadening of the dip are tuned with α. (b) Charging energy U normalized to the coupling broadening Γ as a function of α for different ℓ. Data sets are obtained using a perturbative expression and LSDA for $e^2/\varepsilon l_0 = \hbar\omega_0$.

that while the Kondo regime for the AB interferometer shows a *globally* reduced linear conductance, independently of the gate position,[19,20,27] for the Rashba dot the hallmark of the Kondo effect consists of an *oscillating* $\mathcal{G}_0$ with α, due to the variation of φ.

B. Finite temperature

For finite k_BT the Friedel-Langreth sum rule does not hold and the Green function for the Rashba dot has to be derived explicitly. For that purpose we make use of the equation-of-motion technique.[28] This method is based on differentiating the dot Green function with respect to time generating higher-order Green function's that are approximated following a truncation scheme. Here we choose Lacroix's approximation[28] which goes beyond mean-field theory to include Kondo correlations, evaluating the self-energies in the wide-band limit:

$$G^r_{d\sigma}(\epsilon) = \frac{1-\langle n_{\bar\sigma}\rangle}{\epsilon-\epsilon_0-\Sigma_0-\Sigma^{(1)}_{\rm int}} + \frac{\langle n_{\sigma}\rangle}{\epsilon-\epsilon_0-U-\Sigma_0-\Sigma^{(2)}_{\rm int}}\,, \quad (21)$$

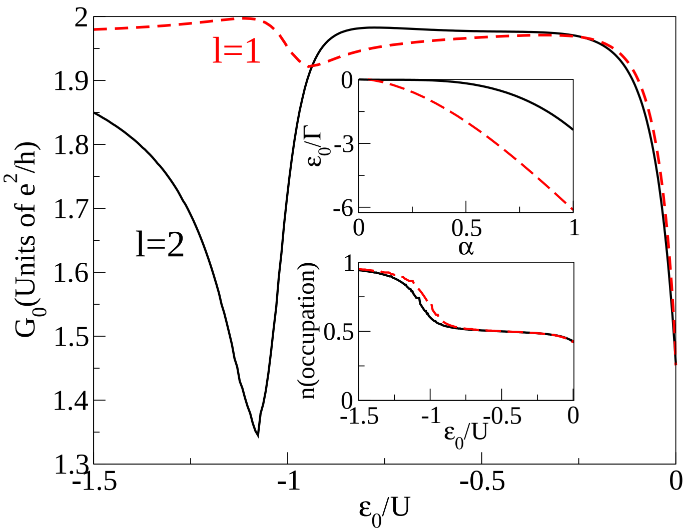

FIG. 2: (Color online) $\mathcal{G}_0$ in the Coulomb-blockade regime for $x = 0.9$, $\ell = 1, 2$ and $k_BT = 0.1\Gamma$ (we set $E_F = 0$). For better comparison, the charging energy remains unchanged with α and ℓ: $U = 7\Gamma$. Upper inset: Nonlinear dependence of ϵ_0 with α. Lower inset: Dot occupation versus gate voltage.

where

$$\Sigma_0 = -i\tilde\Gamma\left(1-ix\cos\varphi\right)\,, \quad (22)$$

$$\Sigma^{(1)}_{\rm int}(\epsilon) = \frac{U\bar\Upsilon(\epsilon)}{\epsilon-\epsilon_0-U-\Sigma_0-\bar\Sigma_1}\,, \quad (23)$$

$$\Sigma^{(2)}_{\rm int}(\epsilon) = -\frac{U\left[\Sigma_1-\bar\Upsilon(\epsilon)\right]}{\epsilon-\epsilon_0-\Sigma_0-\bar\Sigma_1}\,. \quad (24)$$

By defining

$$\Upsilon(\epsilon) = \Psi\left(\frac{1}{2}-i\beta\frac{\epsilon-(2\epsilon_0+U)}{2\pi}\right)-\Psi\left(\frac{1}{2}-i\beta\frac{\epsilon}{2\pi}\right)-i\pi\,, \quad (25)$$

where Ψ is the digamma function, D is bandwidth for the DOS in the leads, and $\beta = 1/k_BT$, we have $\bar\Upsilon(\epsilon) = (\tilde\Gamma/\pi)\Upsilon(\epsilon)\left(1+\pi x\cos\varphi\right)$. Here $\bar\Sigma_1 = -2i\tilde\Gamma\left(1+x\cos\varphi\right)$. Using this result for the Rashba dot Green function in Eq. (9) we can calculate the linear conductance and explore both the Coulomb blockade and the Kondo regimes.

IV. RESULTS

For a noninteracting quantum wire ($U = 0$) with a local Rashba region (see inset in Fig. 1) $\mathcal{G}_0$ versus E_F is shown in Fig. 1(a) for three different values of α. Here, the formation of the Rashba dot can be seen in the appearance of an antiresonance located at the end of the first plateau.[11] Our next question is about the strength of U and how this can affect the previous results. With a finite U the transport through a singly occupied quantum dot is blockaded due to the energy cost required to charge the dot with another electron. To overcome this situation and restore the transport the band bottom of the dot

must be pulled down by applying an external gate. For the present case of a Rashba dot the same effect is accomplished by increasing α. This is the so-called Coulomb blockade regime that arises whenever $U \gg \tilde{\Gamma} \gg k_BT$. At very low temperatures $k_BT \ll k_BT_K$ the Kondo effect occurs because of a strong U in a quantum dot strongly coupled to the leads.[14] In this regime the formation of a singlet state between the itinerant electrons and the localized electron in the quantum dot opens a channel of unitary transmission.[15] In our system we check that these two regimes can be reached for typical values of α. Figure 1(b) displays a calculation of U/Γ obtained from two different approaches. First, we use a perturbative expression,

$$U_{\text{pert}} = -\frac{e^2}{\varepsilon}\int d\vec{r}d\vec{r}'\,\frac{|\phi_0(\vec{r})|^2|\phi_0(\vec{r}')|^2}{|\vec{r}-\vec{r}'|}\,, \qquad (26)$$

where $\phi_0(\vec{r})$ is the wavefunction of the induced bound state by the Rashba interaction and ε is the material permittivity. In this way U is the bare Coulomb interaction which does not include screening effects due to the presence of the rest of charges. To perform a more realistic calculation of U taking into account screening effects we employ the standard local spin-density approximation (LSDA).[29]. In spite of the different approaches used to determine U one clearly sees that both calculations lead to similar results for U. Indeed U is quite large. This means that the effects of a finite U cannot be ignored and, as a result, this changes dramatically the transport properties of this system. Notice that for $\alpha \approx 0$, our calculation for U fails since the wave function in the Rashba dot is now well extended along the quantum wire and not only in the Rashba region. In that case the onsite interaction becomes very weak $U/\tilde{\Gamma} \to 0$ and the noninteracting theory[11] can be safely applied. Hereafter, we consider only situations where the Rashba coupling takes values ranging between $\alpha \approx 0.5 - 1$ [see Fig. 1(b)].

First, we analyze the effect of a finite U that carries this system into the Coulomb blockade regime for $k_BT > k_BT_K$, and the Kondo regime when $-U + E_F + \tilde{\Gamma} < \epsilon_0 < E_F - \tilde{\Gamma}$ and $k_BT \ll k_BT_K$. Typically, in a conventional quantum dot the charging energies for observing correlation effects are in the range of 0.5-1meV with tunneling couplings lying between 100μ eV and 300μ eV.[16] This implies the ratio $U/\Gamma \approx 3 - 10$ which can be easily achieved in our device [see Fig. 1(b)].

Figure 2 displays $\mathcal{G}_0$ versus ϵ_0/U for two different values of ℓ calculated in the Hartree-Fock approximation,[28] where $\Sigma^{(1)}_{\text{int}}(\epsilon) = \Sigma^{(2)}_{\text{int}}(\epsilon) = 0$ in Eq. (21), neglecting in this way correlation effects in the leads. As we mentioned before, α plays the role of the gate voltage and controls the level position ϵ_0/Γ that depends at the same time on ℓ. The upper inset of Fig. 2 shows such dependence for $\ell = 1$ and $\ell = 2$. At these level positions (or their corresponding values of α and ℓ) $\mathcal{G}_0$ consists of two asymmetric Fano resonances located approximately at

$$\begin{aligned}\epsilon_0^{(1)} &\approx -\Sigma_0 = (\tilde{\Gamma}/2)x\cos\varphi\,,\\ \epsilon_0^{(2)} &\approx -U - \Sigma_0 = -U + (\tilde{\Gamma}/2)x\cos\varphi\,.\end{aligned} \qquad (27)$$

These two antiresonances (the resonance at $\epsilon_0 \sim E_F$ is seen in Fig. 2 only partially because $\epsilon_0^{(1)} > 0$ for that value of ϵ_0) arise from the interference between the nonresonant path with the two resonances corresponding to the degenerate points in which the transport through the Rashba dot is energetically allowed. At these points the dot occupation $\langle n_\sigma\rangle$ fluctuates from $0 \to 1/2$ for values of the gate close to $\epsilon_0^{(1)}$ and from $1/2 \to 1$ at values close to $\epsilon_0^{(2)}$, as seen in the lower inset to Fig. 2. Remarkably, the position of the Fano resonances is not very sensitive to a change in ℓ since this only changes the Rashba phase $\varphi = k_\alpha \ell$. We recall that $U \gg \tilde{\Gamma}$ and that, typically, for quantum wires $x \approx 1$ ($T_b \approx 1$). However, we anticipate that in the Kondo regime a change in ℓ modifies dramatically $\mathcal{G}_0$.

Our results for the Kondo regime in the $k_BT = 0$ limit are shown in Fig. 3(a). One might naively think that the "plateau" seen in Fig. 2 between Coulomb-blockade resonances (which amounts to a "valley" in a quantum dot) should decrease due to destructive interference with the nonresonant path, i.e.,

$$\mathcal{G}_0 = \frac{2e^2}{h}(1 - T_b\cos^2\varphi) \approx 0 \quad \text{for } T_b \approx 1\,, \qquad (28)$$

as occurs in the AB interferometer.[19,20] Nevertheless, we see that $\mathcal{G}_0$ versus ϵ_0/U for $\ell = 1$ has a strong dependence on the gate voltage. The dependence is stronger for $\ell = 2$. This is due to the fact that φ is an implicit function of ϵ_0 through α. To confirm this, we plot in Fig. 3(b) $\mathcal{G}_0$ as a function of α for values of $\ell = 1, 2$. Clearly, $\mathcal{G}_0$ in the Kondo plateau is an oscillatory function of α, with the number of oscillations of $\mathcal{G}_0$ being proportional to ℓ. Finally, in Fig. 3(c) we show the total transmission when $\epsilon_0/\Gamma = -1.5$ for different background transmissions using the expression of the Rashba dot Green function given by Eq. (21) that includes Kondo correlations. We see that the Kondo effect arises as a Fano-like resonance pinned at E_F due to its interference with the nonresonant path. As T_b enhances the Kondo dip is shifted towards higher transmission values. This demonstrates that the $\mathcal{T}$ depends not only on U and $\tilde{\Gamma}$ but also on the details of the nonresonant channel via the self-energy.

V. CONCLUSIONS

In short, a quantum wire with a local Rashba field can sustain a Kondo resonance since a quasibound state emerges from the higher subband due to the Rashba interaction. This state is spin degenerate, it strongly couples with the continuum, and our results show that considerable repulsion energy results from charging the dot.

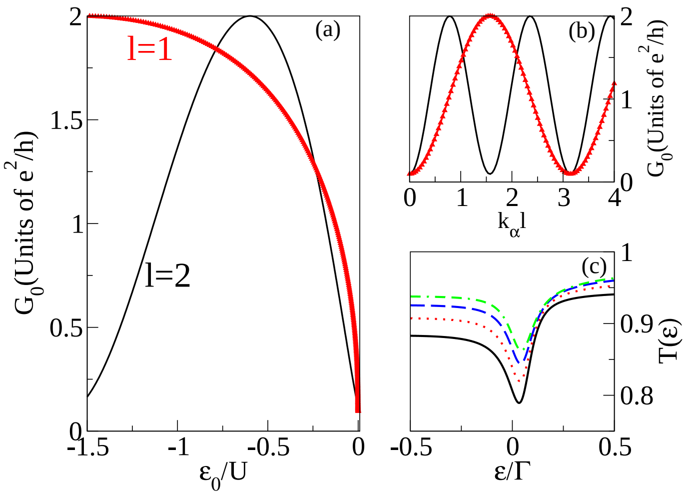

FIG. 3: (Color online) (a) $\mathcal{G}_0$ versus ϵ_0/U in the pure Kondo regime for different dot sizes at $k_BT = 0$ and $x = 0.8$ with $E_F = 0$. (b) $\mathcal{G}_0$ as a function of α. (c) $\mathcal{T}(\epsilon)$ for different background transmissions (from bottom to top: $x = 0.8$, 0.85, 0.9, 0.95) for $\ell = 1$, $\epsilon_0 = -1.5\Gamma$, $U = 7\Gamma$ and $k_BT = 0.05\Gamma$.

The conductance resonances in the Coulomb blockade regime have a Fano form, while in the strong coupling regime we predict an oscillating $\mathcal{G}_0$ as a function of the Rashba strength. We hope that our work sheds light on the relative influence of spin-orbit and electron-electron interactions in nanostructures.

VI. ACKNOWLEDGMENTS

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1. S.A. Wolf, Science **294**, 1488 (2001).
2. D. Loss and D.P. DiVincenzo, Phys. Rev. A **57**, 120 (1998).
3. J. Nitta, T. Akazaki, H. Takayanagi, and T. Enoki, Phys. Rev. Lett. **78**, 1335 (1997). G. Engels, J. Lange, Th. Schäpers, and H. Lüth, Phys. Rev. B **55**, R1958 (1997).
4. E.I Rashba, Fiz. Tverd. Tela (Leningrad) **2**, 1224 (1960). [Sov. Phys. Solid State **2**, 1109 (1960)].
5. M. Khodas, A. Shekter, and A. M. Finkel'stein, Phys. Rev. Lett. **92**, 086602 (2004).
6. M. Valín-Rodríguez, A. Puente, Ll. Serra, Nanotechnology **14**, 882 (2003).
7. L. Zhang, P. Brusheim, and H.Q. Xu, Phys. Rev. B **72**, 045347 (2005).
8. A. Pályi, C. Péterfalvi, and József Cserti, Phys. Rev. B **74**, 073305 (2006).
9. J.C. Egues, G. Burkard, and D. Loss Phys. Rev. Lett. **89**, 176401 (2003).
10. M. Valín-Rodríguez, A. Puente, Ll. Serra, Phys. Rev. B **69**, 153308 (2004).
11. Ll. Serra and D. Sánchez, Phys. Rev. B **74**, 153313 (2006).
12. F. Mireles and G. Kirczenow, Phys. Rev. B **64**, 024426 (2001). M. Governale and U. Zülicke, Phys. Rev. B **66**, 073311 (2002).
13. U. Fano, Phys. Rev. **124**, 1866 (1961).
14. A.C. Hewson, *The Kondo Problem to Heavy Fermions* (Cambridge University Press, Cambridge, UK, 1993).
15. LI. Glazman and M.E. Raikh, Pis'ma Zh. ksp. Teor. Fiz. **47**, 378 (1988) [JETP Lett. **47**, 452 (1988)]. T.K. Ng and P.A. Lee, Phys. Rev. Lett.**61**, 1768 (1988).
16. D. Goldhaber-Gordon, H. Shtrikman, D. Mahalu, D. Abusch-Magder, U. Meirav, and M.A. Kastner, Nature **391**, 156 (1998). S.M. Cronenwett, T.H. Oosterkamp, and L.P. Kouwenhoven, Science **281**, 540 (1998). J. Schmid, J. Weis, K. Eberl, and K.v. Klitzing, Physica B **256-258**, 182 (1998).
17. A. Yacoby, M. Heiblum, D. Mahalu, and Hadas Shtrikman, Phys. Rev. Lett. 74, 4047 (1995).
18. A.L. Yeyati and M. Büttiker Phys. Rev. B **52**, R14360 (1995).
19. B. R. Bulka and P. Stefański Phys. Rev. Lett. **86**, 5128 (2001).
20. W. Hofstetter, J. König, and H. Schoeller Phys. Rev. Lett. **87**, 156803 (2001).
21. Q.F. Sun, J. Wang, and H. Guo, Phys. Rev. B **71**, 165310 (2005).
22. W. Rudzinski and J. Barnas, Phys. Rev. B **64**, 085318 (2001); R. López and D. Sánchez, Phys. Rev. Lett. **90**, 116602 (2003); M.-S. Choi, D. Sánchez and R. López, Phys. Rev. Lett. **92**, 056601 (2004); B. Dong, G.H. Ding, H.L. Cui, X.L. Lei, Europhys. Lett. **69**, 424 (2005).
23. K.C. Lin and D.S. Chuu Phys. Rev. B **72**, 125314 (2005).
24. K. Kobayashi, H. Aikawa, S. Katsumoto, and Y. Iye Phys. Rev. Lett. **88**, 256806 (2002).
25. D.Sánchez, Ll. Serra, and R. López, (unpublished).
26. D.C. Langreth Phys. Rev. 150, 516 (1966).
27. M. Sato, H. Aikawa, K. Kobayashi, S. Katsumoto, and Y. Iye Phys. Rev. Lett. **95**, 066801 (2005).
28. C. Lacroix, J. Phys. F **11**, 2389 (1981). Y. Meir, N. S. Wingreen, and P. A. Lee Phys. Rev. Lett. **66**, 3048 (1991).
29. S. M. Reimann and M. Manninen, Rev. Mod. Phys. **74**, 1283 (2002).